

An Unusual Two Higgs Doublet Model from Warped Space

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Abstract

We study a simple two Higgs doublet model (2HDM) in the Randall-Sundrum scenario, with an IR brane localized Higgs field and a second doublet arising from a $t\bar{t}$ condensate due to strong Kaluza-Klein gluon effects. The effective 2HDM predicts that the ratio of the brane to condensate vacuum expectation values $\tan\beta \sim 3$. It also predicts a standard model like Higgs boson of mass $\mathcal{O}(100)$ GeV and a heavier scalar at the scale of the lowest KK gluon mass, which we take to be $M_{KK} \gtrsim 1.5$ TeV. The pseudoscalar and the charged scalars are degenerate in mass at tree-level and are $\mathcal{O}(M_{KK})$. There are no tree-level flavor changing neutral currents (FCNC) for the down-type quarks and the standard model results hold there. In contrast, FCNC decays of the t-quark larger than in the SM are expected.

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I. INTRODUCTION

With the successful start up of the Large Hadron Collider (LHC) we can look forward to the discovery of the electroweak symmetry breaking mechanism. While the Standard Model (SM) with only one Higgs doublet is highly successful phenomenologically, it is theoretically unsatisfactory due to the Higgs boson mass being quadratically sensitive to unknown physics at high energies. This is the hierarchy problem and is the main motivation for many models that extend the SM. Common to many models is an extended Higgs sector. For example, the supersymmetric extension of the SM which addresses the hierarchy problem requires two Higgs doublets, distinguished by their opposite $U(1)_Y$ charges.

Another proposed solution to the hierarchy problem is the introduction of an extra dimension with a warp factor [1]. These constructions, referred to as Randall-Sundrum (RS) models, are also interesting flavor models when fermions become bulk fields [2]. The observed fermion mass hierarchy is given by the same warp factor which determines the electroweak-Planck hierarchy: the mass matrices are determined by fermion locations in the bulk, without fine tuning the Yukawa couplings. Generically the Kaluza-Klein (KK) excitations have masses in the TeV range and above in order to satisfy the many electroweak precision measurements. The lower limit of ~ 2 to 3 TeV for KK gauge bosons make these models testable at the LHC.

Many recent studies within this RS framework of flavor reveal that the custodial symmetry of the SM Higgs potential that protects the T parameter from large corrections is not automatic in 5D warped models [3]. An elegant solution is to add an $SU(2)_R$ custodial gauge symmetry which maintains tree level protection of T [4]. With an additional discrete left-right symmetry P_{LR} the Zb_Lb_L coupling will also be protected from large corrections [5]. Moreover, sizable deviations of the Ztt and Wtb couplings from the SM are expected and will be something to search for at the LHC.

It is a generic feature of RS flavor models that the t-quark, being much heavier than the other fermions, is located near the TeV (or IR) brane. The KK gauge bosons also have profiles peaking close to the IR brane. In particular, the KK gluons, G_{KK} , are expected to have large wavefunction overlaps with t-quarks of both handedness. Thus in terms of the 4D effective theory the interactions $G_{KK}t\bar{t}$ are of strong coupling strength. The exact magnitude depends on the parameters c_L^3, c_R^3 that characterize the location of t_L and t_R in warped space.

At low energies, i.e. below M_{KK} (the mass of the first G_{KK}) we might expect the t-quarks to form Nambu-Jona-Lasinio condensates [6]. Specifically we examine the condensing of the SM third family doublet with t_R via G_{KK} exchanges to form a composite Higgs doublet. We do not expect the lighter quarks to similarly condense since they are located further away from the IR brane, and hence the effective couplings with the G_{KK} are much smaller. This observation is not new and has been exploited to generate the Planck scale in [7]. A similar condensate of a fourth generation has also been studied in [8]. Here we study the physics of this composite doublet using the a minimal custodial RS framework described in Ref. [4]. For a review of RS models see [9]. Ref.[10] gives a thorough review of top quark condensates physics.

In this paper we also assume that there is the usual SM Higgs doublet, H , which is localized on the IR brane.¹ It has the usual Yukawa couplings to the SM fermions. Below M_{KK} the effective theory has both an ‘elementary’ and a composite Higgs, Φ . Electroweak symmetry breaking will involve both these scalars. Since Φ only couples to t_R there are no flavor changing neutral currents (FCNC) in the down type quark sector. We do expect FCNC involving the t quarks to be important. The t-quark mass receives contributions from both the elementary and the composite Higgs, so one can move the localization of Q_3 away from the IR brane and still get a good agreement with the experimental value. This ameliorates the Zb_Lb_L problem and we find that the extra P_{LR} is not always required.

Another interesting feature of the model is the simplicity of the effective scalar potential $V(\Phi, H)$. Whereas a generic 2HDM potential contains fourteen real parameters, our model contains only eight (four of which only appear in one of two specific combinations). Further, all of these parameters occur already in the RS framework of Ref. [4]; they include the AdS curvature k , the 5D size r_c (or equivalently the mass M_{KK}), the two c parameters, and the 5D Yukawa couplings of the elementary scalar. The relative simplicity of the scalar sector makes our model quite predictive; for example, the pseudoscalar and the charged Higgs are degenerate in mass at tree level. We also find that this extended scalar sector automatically respects the $U(1)_{em}$ and CP symmetries.

In section II we describe the condensate formation using G_{KK} and construct $V(\Phi, H)$

¹ A similar idea has previously been studied in Ref. [11, 12]; however, both models lacked an explicit UV completion of the 2HDM, and further in [11] the top quark only coupled to the composite Higgs.

which is valid below M_{KK} . In Sec. III, we study the pattern of electroweak symmetry breaking and the spectrum of scalar states. Numerical results are given in Sec. IV, and a discussion of flavor-changing neutral currents in Sec. V. Finally we give our conclusions in Sec. VI.

II. CONDENSATE FROM WARP SPACE

The model we study is the minimal custodial RS model (MCRS) given in [4]. Here we give a brief description of the model following the notations of [13]. The bulk gauge symmetry is $SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_X$. The background geometry is a slice of AdS_5 space specified by the metric

$$ds^2 = G_{AB} dx^A dx^B = e^{-2\sigma(\phi)} \eta_{\mu\nu} dx^\mu dx^\nu - r_c^2 d\phi^2, \quad (1)$$

where $\sigma(\phi) = kr_c|\phi|$, $\eta_{\mu\nu} = \text{diag}(1, -1, -1, -1)$, k is the AdS_5 curvature, and $-\pi \leq \phi \leq \pi$. The theory is compactified on an $S_1/(\mathbb{Z}_2 \times \mathbb{Z}'_2)$ orbifold, with r_c the radius of the compactified fifth dimension, and the orbifold fixed points at $\phi = 0$ and $\phi = \pi$ correspond to the UV (Planck) and IR (TeV) branes respectively. To solve the hierarchy problem, $k\pi r_c$ is set to ≈ 37 . The warped down scale is defined to be $\tilde{k} = k e^{-k\pi r_c}$. Note that \tilde{k} sets the scale of the first KK gauge boson mass, $m_{gauge}^{(1)} \approx 2.45\tilde{k}$, which determines the scale of the new KK physics.

The SM quarks are embedded into $SU(2)_L \times SU(2)_R \times U(1)_X$ via the five-dimensional (5D) bulk Dirac spinors

$$Q_i = \begin{pmatrix} u_{iL} [+ , +] \\ d_{iL} [+ , +] \end{pmatrix}, \quad U_i = \begin{pmatrix} u_{iR} [+ , +] \\ \tilde{d}_{iR} [- , +] \end{pmatrix}, \quad D_i = \begin{pmatrix} \tilde{u}_{iR} [- , +] \\ d_{iR} [+ , +] \end{pmatrix}, \quad i = 1, 2, 3, \quad (2)$$

where Q_i transforms as $(2, 1)_{1/6}$, and U_i, D_i transform as $(1, 2)_{1/6}$. The parity assignment \pm denotes the boundary conditions applied to the spinors on the [UV, IR] brane, with $+$ ($-$) being the Neumann (Dirichlet) boundary conditions. Only fields with the $[+, +]$ parity contain zero-modes that do not vanish on the brane. These survive in the low energy spectrum of the 4D effective theory, and are identified as the SM fields.

The profiles of the zero-mode fermions are given by their flavor functions

$$f_{L,R}^0(\phi, c_{L,R}) = \sqrt{\frac{k\pi r_c (1 \mp 2c_{L,R})}{e^{k\pi r_c (1 \mp 2c_{L,R})} - 1}} e^{(1/2 \mp c_{L,R})kr_c \phi}, \quad (3)$$

where the c -parameter is determined by the bulk Dirac mass parameter, $m = ck$, and the upper (lower) sign applies to the LH (RH) label. One of the fermion chiralities is projected out of the zero-mode; which one depends on the fermion's orbifold parity.

With the introduction of an IR brane localized Higgs field H the Yukawa interactions are all localized there. After spontaneous electroweak symmetry breaking the contribution to the mass matrices of the SM fermions are

$$(M_f^{RS})_{ij} = \frac{v_H}{\sqrt{2}k\pi r_c} y_{5,ij}^f f_L^0(\pi, c_{f_i}^L) f_R^0(\pi, c_{f_j}^R) \equiv \frac{v_H}{\sqrt{2}k\pi r_c} y_{5,ij}^f F_L(c_{f_i}^L) F_R(c_{f_j}^R), \quad f = u, d, \quad (4)$$

where the label f denotes up-type or down-type quark species. $y_{5,ij}^f$ denotes the 5D Yukawa couplings, ij are family symbols, and v_H is the VeV of H .

In the gauge (weak) eigenbasis, the coupling of the n th level KK gluon, $G^{(n)}$, to zero-mode fermions is given by

$$G_\mu^{A(n)} \left[\sum_i (g_f^n)_{ii}^L \bar{f}_{iL} T^A \gamma^\mu f_{iL} + (L \rightarrow R) \right], \quad f = u, d, \quad (5)$$

where i is a generation index, T^A are the generators of $SU(3)$, and $(g_f^n)_{ii} = \text{diag}(g_{f_1}^n, g_{f_2}^n, g_{f_3}^n)$ is the weak eigenbasis coupling matrix

$$g_{f_i}^n = \frac{g_s}{\pi} \int_0^\pi d\phi |f^0(\phi, c_{f_i})|^2 \chi_n(\phi), \quad g_s = \frac{g_{5s}}{\sqrt{r_c \pi}}. \quad (6)$$

Here, g_{5s} is the bulk 5D $SU(3)$ gauge coupling, g_s that in the SM, and χ_n the profile of the n th KK gluon. We will exclusively consider the coupling of the G_{KK}^1 to the t-quarks and thus write

$$g_L = \frac{g_s}{\pi} \int_0^\pi d\phi |f^0(\phi, c_{f_{3L}})|^2 \chi_1(\phi) \quad (7a)$$

$$g_R = \frac{g_s}{\pi} \int_0^\pi d\phi |f^0(\phi, c_{f_{3R}})|^2 \chi_1(\phi). \quad (7b)$$

For processes with small external momenta, tree-level exchange of G_{KK}^1 leads to 4-Fermi interactions between zero mode fermions given by

$$-\frac{g_i g_j}{M_{KK}^2} (\overline{Q_{iL}} T^A \gamma^\mu Q_{iL}) (\overline{f_{jR}} T^A \gamma_\mu f_{jR}) = \frac{g_i g_j}{M_{KK}^2} (\overline{Q_{iL}} f_{jR}) (\overline{f_{jR}} Q_{iL}) + O(1/N_c) \quad (8)$$

where $N_c = 3$ for $SU(3)_c$. Taking $Q_i = Q_3$ and $f_j = t_R$, this generates the t-quark condensate considered in [6].

The exchange of spin-1 KK-gluons leads to an attractive force in Eq. (8). A condensate can only form if the effective couplings' g_i 's are sufficiently strong. In the RS model this requires the overlap of the KK gluons and the fermion zero modes be large and thus enhance the strong coupling g_s . The wavefunction of the lightest KK gluon G_{KK}^1 peaks near the IR brane; so the only fermions that can condense are those which are similarly localized. Eq. (4) shows that light fermions are oriented towards the UV brane, and so no enhancement of g_s is expected. The large top mass requires the fields Q_3 and t_R to be IR localized; as we shall see later, this remains true in the presence of the composite Higgs, since a substantial part of the t-quark mass still comes from H . This justifies our restriction of Eq. (8) to only Q_3 and t_R .

We have focused on the role of G_{KK}^1 in condensate formation. The higher KK-gluon modes contribute only a small correction, due both to a smaller overlap integral from Eq. (6) and their higher masses. Numerical calculations show that the contribution of these states to the coefficient in Eq. (8) is at most $\sim 5\%$ for parameters of interest (and usually much less). Similarly, the KK fermions give small corrections to the above picture when they are integrated out.

Below the scale of M_{KK} the condensate in Eq. (8) can be viewed as a composite Higgs doublet denoted by Φ [6]. It has the same $SU(2)_L \times U(1)_Y$ quantum numbers as the SM Higgs. This is in addition to the elementary scalar field H which is in the MCRS model. At the scale M_{KK} , Φ is a static auxiliary field. Following [14] we can write the Lagrangian at scale M_{KK} as

$$\mathcal{L} = |D_\mu H|^2 - m_0^2 H^\dagger H - \frac{1}{2} \lambda_0 (H^\dagger H)^2 + \lambda_t \overline{Q_L} t_R \tilde{H} + g_t \overline{Q_L} t_R \tilde{\Phi} - M_{KK}^2 \Phi^\dagger \Phi + h.c. \quad (9)$$

where $\tilde{H} = i\sigma_2 H^*$, $\tilde{\Phi} = i\sigma_2 \Phi^*$, $Q_L = (t, b)_L$, and m_0^2 , λ_0 are the parameters in the brane Higgs scalar potential. We have omitted the Yukawa interactions of the brane Higgs with light fermions as they play no role here. D_μ denotes the gauge covariant derivative. Integrating out Φ reproduces the 4-Fermi interaction of Eq. (8). For the Yukawa couplings, $g_t \equiv \sqrt{g_L g_R}$ is given in Eq. (7), and λ_t is determined by the wavefunction overlap of Q_3 , t_R , and H at the IR brane. Specifically, using the notation of Eq. (4) and defining $y_{5D} \equiv y_{5,33}^u$ we have

$$\lambda_t = \frac{y_{5D}}{k\pi r_c} F_L(c_{3L}) F_R(c_{tR}) = y_{5D} \sqrt{\frac{(1 - 2c_{3L})(1 + 2c_{tR})}{(1 - e^{k\pi r_c(2c_{3L}-1)})(1 - e^{-k\pi r_c(2c_{tR}+1)})}}. \quad (10)$$

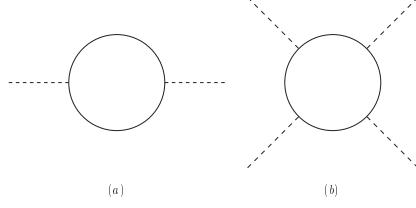


FIG. 1: Fermion bubble contribution to scalar (a) 2-point functions and (b) 4-point functions. The dashed lines can be either Φ or H fields.

At scales $\mu < M_{KK}$, quantum fluctuations generate a kinetic term for Φ as well as kinetic and mass term mixings between ϕ and H . We will calculate these corrections in the fermion bubble approximation, depicted in Fig. 1. The effective Lagrangian takes the form

$$\begin{aligned}
\mathcal{L} = & [1 + \lambda_t^2 \epsilon] |D_\mu H|^2 + \lambda_t g_t \epsilon [(D_\mu H)^\dagger D^\mu \Phi + h.c.] + g_t^2 \epsilon |D_\mu \Phi|^2 \\
& - [m_0^2 - \lambda_t^2 \Delta^2] H^\dagger H + \lambda_t g_t \Delta^2 [H^\dagger \Phi + \Phi^\dagger H] - [M_{KK}^2 - g_t^2 \Delta^2] \Phi^\dagger \Phi \\
& - \left[\frac{1}{2} \lambda_0 + \lambda_t^4 \epsilon \right] (H^\dagger H)^2 - 2 \lambda_t^2 g_t^2 \epsilon \left[\frac{1}{2} (H^\dagger \Phi + \Phi^\dagger H)^2 + H^\dagger H \Phi^\dagger \Phi \right] \\
& - 2 \lambda_t^3 g_t \epsilon H^\dagger H (H^\dagger \Phi + \Phi^\dagger H) - 2 \lambda_t g_t^3 \epsilon \Phi^\dagger \Phi (H^\dagger \Phi + \Phi^\dagger H) - g_t^4 \epsilon (\Phi^\dagger \Phi)^2 \\
& + \lambda_t \overline{Q}_L t_R \tilde{H} + g_t \overline{Q}_L t_R \tilde{\Phi} + h.c.
\end{aligned} \tag{11}$$

Here,

$$\begin{aligned}
\epsilon &= \frac{N_c}{16\pi^2} \ln \left(\frac{M_{KK}^2}{\mu^2} \right); \\
\Delta^2 &= \frac{2N_c}{16\pi^2} (M_{KK}^2 - \mu^2),
\end{aligned} \tag{12}$$

are calculated in the 1-loop approximation. We have also taken the cutoff to be M_{KK} , above which the 4-Fermi condensate approximation is no longer valid.

The kinetic mixing term can be diagonalized in the customary fashion. The transformations

$$\begin{aligned}
H &= \hat{H} \\
\Phi &= -\frac{\lambda_t}{g_t} \hat{H} + \frac{1}{g_t \sqrt{\epsilon}} \hat{\Phi}
\end{aligned} \tag{13}$$

will cast the kinetic terms into canonical diagonalized form. The resulting Lagrangian of the scalars is delightfully simple:

$$\mathcal{L} \supset |D_\mu \hat{H}|^2 + |D_\mu \hat{\Phi}|^2 - V(\hat{H}, \hat{\Phi})$$

where the scalar potential is given by

$$V(\hat{H}, \hat{\Phi}) = \left(m_0^2 + \frac{\lambda_t^2}{g_t^2} M_{KK}^2 \right) \hat{H}^\dagger \hat{H} - \frac{\lambda_t}{g_t^2 \sqrt{\epsilon}} M_{KK}^2 \left(\hat{H}^\dagger \hat{\Phi} + \hat{\Phi}^\dagger \hat{H} \right) + \left(\frac{M_{KK}^2}{g_t^2 \epsilon} - \frac{\Delta^2}{\epsilon} \right) \hat{\Phi}^\dagger \hat{\Phi} + \frac{1}{2} \lambda_0 (\hat{H}^\dagger \hat{H})^2 + \frac{1}{\epsilon} (\hat{\Phi}^\dagger \hat{\Phi})^2 \quad (14)$$

In this basis the only interaction between the two scalar doublets is the mass mixing. This mixing is related to the c parameters via g_t and λ_t . We reiterate that this potential is valid in the range $\mu < M_{KK}$.

III. ELECTROWEAK SYMMETRY BREAKING OF 2HDM

The particular form of our Higgs potential makes study of the scalar VEVs and physical states straightforward. It is notable that in this model, alignment of the Higgs VEVs is automatic from the structure of the potential; while CP is conserved since all parameters are real. We denote the vacuum expectation values of \hat{H} and $\hat{\Phi}$ by v_H and v_ϕ respectively, and also define $\tan \beta = \frac{v_H}{v_\phi}$. Minimizing $V(\hat{H}, \hat{\Phi})$ yields two coupled cubic equations

$$\left(m_0^2 + \frac{\lambda_t^2}{g_t^2} M_{KK}^2 \right) v_H - \frac{\lambda_t}{g_t^2 \sqrt{\epsilon}} M_{KK}^2 v_\phi + \frac{\lambda_0}{2} |v_H|^2 v_H = 0, \quad (15)$$

and

$$\left(\frac{M_{KK}^2}{g_t^2 \epsilon} - \frac{\Delta^2}{\epsilon} \right) v_\phi - \frac{\lambda_t}{g_t^2 \sqrt{\epsilon}} M_{KK}^2 v_H + \frac{2}{\epsilon} |v_\phi|^2 v_\phi = 0. \quad (16)$$

We require $v_H^2 + v_\phi^2 = v^2$ where $v = 246$ GeV.

The physical scalar spectrum consists of a pair of charged Higgs H^\pm , a pseudoscalar A , and two neutral scalars $h_{1,2}$. The charged and pseudoscalar sectors have the same mass matrix

$$M_\pm^2 = \begin{pmatrix} a + \frac{\lambda_0}{2} v_H^2 & c \\ c & b + \frac{1}{\epsilon} v_\phi^2 \end{pmatrix}, \quad (17)$$

where

$$a = m_0^2 + \frac{\lambda_t^2}{g_t^2} M_{KK}^2; \quad (18)$$

$$b = \frac{1}{\epsilon} \left(\frac{M_{KK}^2}{g_t^2} - \Delta^2 \right); \quad (19)$$

$$c = -\frac{\lambda_t}{g_t^2 \sqrt{\epsilon}} M_{KK}^2. \quad (20)$$

Since $c^2 = (a + \frac{\lambda_0}{2}v_H^2)(b + \frac{1}{\epsilon}v_\phi^2)$, this matrix has a null eigenvalue, corresponding to the eaten Goldstone modes. The non-zero eigenvalue gives the mass of the physical states, and hence at tree level we have

$$M_A^2 = M_H^2 = \frac{2\lambda_t}{g_t^2\sqrt{\epsilon}\sin 2\beta}M_{KK}^2 \quad (21)$$

We shall see in section IV that these masses are expected to be $\lesssim M_{KK}^2$.

The degeneracy of the charged and pseudoscalar Higgses can be understood from a symmetry stand point. The 2HDM has the maximal symmetry of $SO(8)$. Without the mass mixing term the $\hat{\Phi}$ and \hat{H} will have the individual symmetries $SU(2)_{\hat{\Phi}_L} \times SU(2)_{\hat{\Phi}_R}$ and $SU(2)_{\hat{H}_L} \times SU(2)_{\hat{H}_R}$, and their cross product is a subgroup of $SO(8)$. Spontaneous symmetry breaking yields

$$\begin{aligned} SU(2)_{\hat{\Phi}_L} \times SU(2)_{\hat{\Phi}_R} &\xrightarrow{v_\phi} SU(2)_{D\hat{\Phi}}; \\ SU(2)_{\hat{H}_L} \times SU(2)_{\hat{H}_R} &\xrightarrow{v_H} SU(2)_{D\hat{H}}, \end{aligned} \quad (22)$$

and the $SU(2)_D$ are analogous to the custodial $SU(2)$ of the SM scalar potential after symmetry breaking. These symmetries are then softly broken by the mixing term in $V(\hat{\Phi}, \hat{H})$, which indicates that there is a remaining custodial symmetry $SU(2)_V \subset SU(2)_{D\hat{\Phi}} \times SU(2)_{D\hat{H}}$. The three Goldstones form a triplet under this $SU(2)_V$. The two charged Higgs and the pseudoscalar form another triplet. The remaining two scalars are singlets and complete the number of fields. Hence, at tree level we have $M_A = M_H$.

The mass squared matrix of the two scalars is given by

$$M_0^2 = \begin{pmatrix} a + \frac{3}{2}\lambda_0 v_H^2 & c \\ c & b + \frac{3}{\epsilon}v_\phi^2 \end{pmatrix} \quad (23)$$

Taking the trace and the determinant of this matrix gives $\text{Tr } M_0^2 = M_H^2 + k_1 v^2$ and $\det M_0^2 = k_2 M_H^2 v^2$, where $k_{1,2}$ are ratios of $\mathcal{O}(1)$ parameters. This implies that one of the scalars has a mass $\sim M_H \sim \mathcal{O}(\text{TeV})$, while the other has mass of order of electroweak scale. Heavy scalars are expected in this model. This can be seen from the scalar potential given in Eq. (14), since the quartic term $(\hat{\Phi}^\dagger \hat{\Phi})^2$ has a large coefficient. We close this section by giving the mixing angle α between the two scalars:

$$\tan 2\alpha = -\frac{2c}{a + \frac{3}{2}\lambda_0 v_H^2 - b - \frac{3}{\epsilon}v_\phi^2}, \quad (24)$$

and

$$\begin{pmatrix} H_0 \\ h_0 \end{pmatrix} = \begin{pmatrix} \cos \alpha & -\sin \alpha \\ \sin \alpha & \cos \alpha \end{pmatrix} \begin{pmatrix} \hat{H}_R \\ \hat{\Phi}_R \end{pmatrix}. \quad (25)$$

The subscript R stands for the real part expansion around the VeVs, and h_0 (H_0) is the physical lighter (heavier) neutral scalar. We will show in section IV that the couplings of h_0 to gauge bosons are almost identical to those of the SM Higgs; we are in the ‘decoupling’ limit of 2HDMs [15].

IV. T-QUARK MASS AND NUMERICAL ANALYSIS

We start with a discussion of the top quark mass generation. Below the cut-off scale of M_{KK} the 2HDM is valid. After making the field rotation of Eq. (13), the Yukawa sector for the 3rd generation becomes

$$\mathcal{L}_Y = \lambda_t \overline{Q}_L t_R \tilde{H} + g_t \overline{Q}_L t_R \tilde{\Phi} + h.c. \rightarrow \frac{1}{\sqrt{\epsilon}} \overline{Q}_L t_R \tilde{\Phi} + h.c. \quad (26)$$

Eq. (26) shows that the t-quark gets its mass from coupling to $\hat{\Phi}$, which after symmetry breaking gives²

$$m_t = \frac{v \cos \beta}{\sqrt{2\epsilon}}. \quad (27)$$

We have already assumed that v has the SM value; Eq. (27) then tells us that the top quark mass determines the value of β . This determination of *both* Higgs field VeVs is unusual in 2HDMs.

To make use of this result, we need to specify the renormalization scale μ . It is reasonable to choose $\mu = v$. Using the value of m_t at the electroweak scale fixes $\cos \beta \sim \sqrt{\epsilon}$. Eq. (16) then becomes a constraint on λ_t , g_t , with the allowed regions shown in Fig. 2. Recall that $g_t = \sqrt{g_L g_R}$ is related to the parameters c_L^3 , c_R^3 via Eq. (7), and λ_t is related to them via Eq. (10); so for a given y_{5D} , this becomes a restriction on the 5D mass parameters of the third generation.

Taking into account the above considerations, Eq. (15) now constrains the m_0^2 and λ_0 parameters of the brane Higgs; see Fig. 3 for the results. We see that while m_0 is of the

² We are neglecting the contributions from Yukawas involving H that mix generations, since these terms will be small.

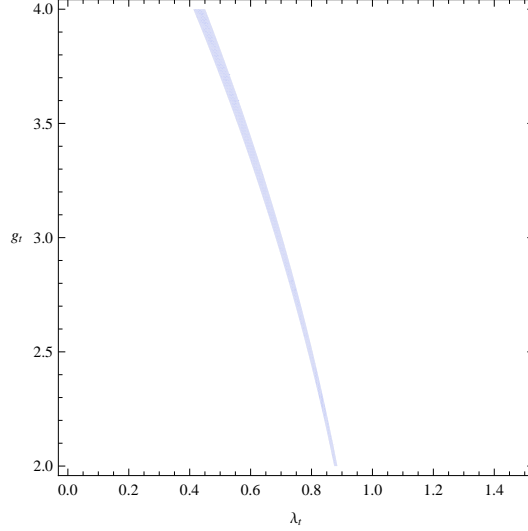


FIG. 2: Allowed region in the $\{\lambda_t, g_t\}$ parameter space that satisfies Eqs. (27) and (16). Note that for solutions in the region shown, we must take $\cos \beta$ positive. The spread corresponds to taking M_{KK} to lie between 1.5 to 4 TeV and experimentally allowed region of m_t from 169.7 to 172.9 GeV.

order of the weak scale the quartic coupling of the brane Higgs field is of $\mathcal{O}(1)$ and not smaller; but the coupling remains well within the perturbative regime. We also see that, after imposing the various constraints, there remain three free parameters in our 2HDM, which we take to be g_t , λ_0 and M_{KK} .

Numerically solving for the spectrum of the scalars, we predict the mass of the lighter Higgs boson to be in the range of 100–400 GeV; while the heavier scalar, pseudoscalar and charged Higgs are all near $M_{KK} \sim \text{TeV}$. We plot the lightest scalar mass in Fig. 4 and the others in Fig. 5.

For h_0 , the mass is determined almost entirely by λ_0 , with essentially no dependence on g_t . As Fig. 4 demonstrates, varying M_{KK} also has very little effect on the mass of this scalar. We further note that this Higgs boson has essentially the properties of the SM Higgs; the mixing angles α and β are almost equal, with the choice of 2α in the second quadrant (see Eq. (24)). All these properties can be qualitatively understood from the mass squared matrices of the neutral and charged scalars, Eq. (17) and Eq. (23). We can write

$$M_0^2 = M_{\pm}^2 + \begin{pmatrix} \lambda_0 v^2 \sin \beta & 0 \\ 0 & 4m_t^2 \end{pmatrix}. \quad (28)$$

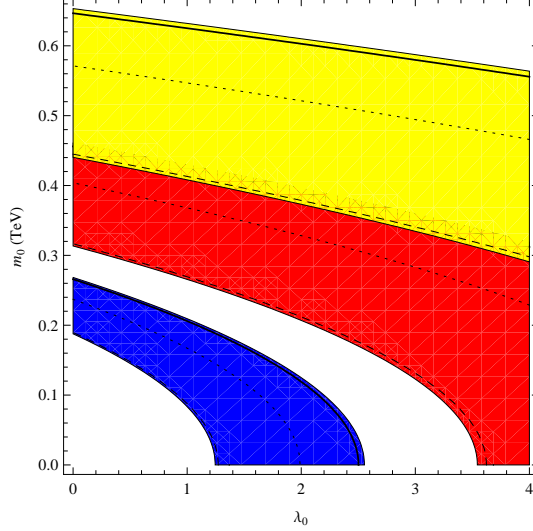


FIG. 3: Allowed region in the $\{\lambda_0, m_0\}$ parameter space that satisfies Eq. (15). The blue, red and yellow regions correspond to $M_{KK} = 1.5, 2.5$ and 3.5 TeV respectively. The different lines (solid, dotted, dash) correspond to different values of $g_t = 2, 3, 4$. We have taken $\sin \beta$ positive, which implies m_0^2 positive, and the central experimental value for m_t .

The charged Higgs masses are $\mathcal{O}(M_{KK})$ and the matrix M_{\pm}^2 is diagonalized by a rotation matrix with angle β . The second term of Eq. (28) can then be treated as a perturbation, being of $\mathcal{O}(.01)$. This implies $\beta \simeq \alpha$, that one of the two scalars should be nearly degenerate with the charged/pseudoscalar Higgses, and that the other should be at the electroweak scale. Using Eq. (28) we find an approximate expression for the mass of h_0 :

$$M_{h_0}^2 \simeq \lambda_0 v^2 \sin^4 \beta + 2\epsilon m_t^2. \quad (29)$$

The near equality of α and β implies that the lighter scalar is very SM like. For example, its coupling to the Z boson has a factor of $\cos(\beta - \alpha)$ relative to the SM coupling; so it has essentially SM strength. The LEP bound on the Higgs mass of > 114 GeV [16], as well as the Tevatron bound $m \notin (163, 166)$ GeV [17], apply and are shown in Fig. 4. A first glance at Eq. (26) would suggest that the light scalar would have no coupling to t-quarks. However, we will show in section V that the couplings of h_0 to all fermions are proportional to $m_f \sin \alpha / \sin \beta \simeq m_f$. We conclude that the lighter scalar is indistinguishable from the SM Higgs.

For the heavier scalars, the masses are mostly determined by g_t , exhibiting little dependence on λ_0 . This is congruent with our expectations from Eq. (21) and Eq. (28). Note that

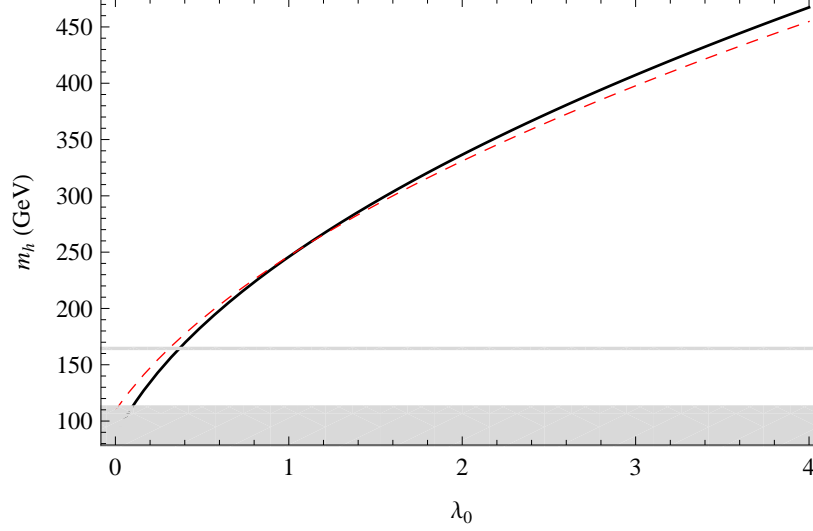


FIG. 4: Value of the mass of the lighter Higgs boson, as a function of λ_0 . The black solid line is for $M_{KK} = 1.5$ TeV and the red dashed line is for $M_{KK} = 4$ TeV. The shaded regions are the LEP and Tevatron exclusions for the Higgs mass.

as shown in Fig. 5, all these states are nearly degenerate with the scalar slightly heavier than the charged and the pseudoscalar states. The splittings are of $\mathcal{O}(10)$ GeV. We also note that the heavier scalar coupling to the Z-boson is negligible, since it is proportional to $\sin(\beta - \alpha)$. This state can be produced at the LHC if kinematically allowed, but will not be produced in a e^+e^- collider in the ZH channel.

An alternative way of obtaining m_t is via the gap equation with a brane Higgs contribution. Instead of Eq.(27) we have

$$\begin{aligned}
 m_t &= \frac{\lambda_t}{\sqrt{2}} v \sin \beta - i \frac{N_c g_L g_R}{2M_{KK}^2} \int \frac{d^4 \ell}{(2\pi)^4} \frac{\text{Tr}(\ell + m)}{\ell^2 - m_t^2} \\
 &= \frac{\lambda_t}{\sqrt{2}} v \sin \beta + \frac{N_c g_L g_R m_t}{8\pi^2} \left[1 + \frac{m_t^2}{M_{KK}^2} \ln \frac{m_t^2}{M_{KK}^2} \right]
 \end{aligned} \tag{30}$$

The graphical depiction of this is given in Fig. 6. This approach includes non-perturbative effects via the fermion loop. The range of masses predicted for h_0 is between 110 to 470 GeV which is very close to that from the 2HDM (see Fig. 4). We find that the approaches agree with each other to within twenty percent in the observables we calculated. This is consistent with the approximation of dropping $1/N_c$ terms. Fig. 7 is an example of the qualitative agreement and quantitative difference of the two calculations.

Although it is difficult to quantify at what coupling strength fermion condensates will

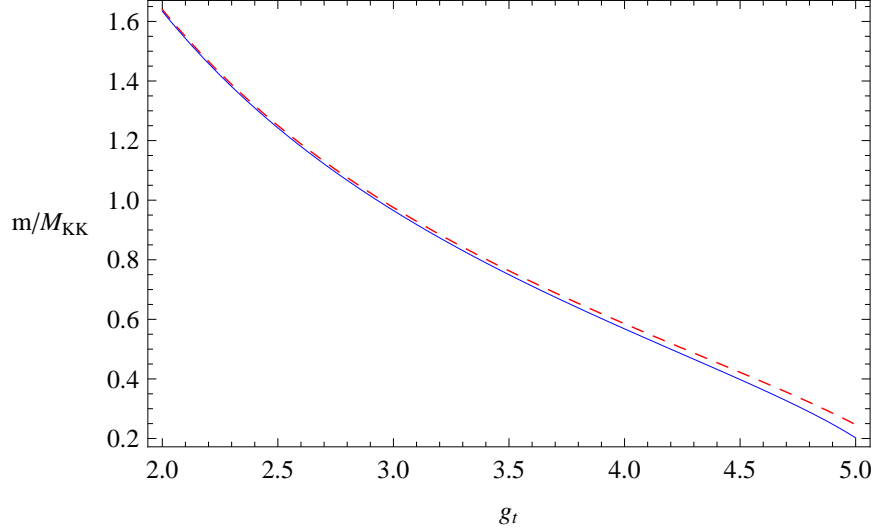


FIG. 5: The masses of the charged, pseudoscalar and heavier scalar Higgses, as a function of g_t and relative to M_{KK} . The charged and pseudoscalar states are degenerate at tree level and are shown by the blue, solid line; the heavier scalar state is the red, dashed line.

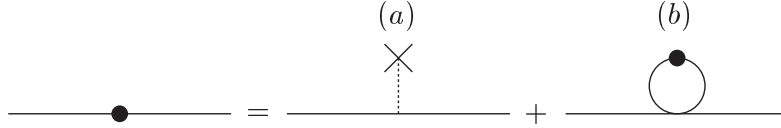


FIG. 6: The gap equation for m_t . (a) Contribution from the brane Higgs (dash line) and the cross denotes the VeV. (b) The fermion bubble contribution.

form, we can draw from our experience with low energy QCD. We expect condensates to form if the effective coupling $g_t \gtrsim g_s(2 \text{ GeV}) \sim 2.1$. Our solutions shown in Fig. 2 are consistent with that. In Fig. 8 we show how the effective condensate coupling g_t is related to the c -parameters. In order to satisfy both the EWPT and the constraints from the correction to the SM $\delta g_b^L < 0.01$ we require $c_L^3 > 0.3$ [4, 13] and this is marked in Fig. 8. In Fig. 9 we show the allowed region and the contour of 5D Yukawa couplings that gives the t-quark mass. Hence, we conclude that, with the aid of a t-quark condensate, the parameter space for the minimal custodial RS model does not require a further discrete symmetry to satisfy all the phenomenological constraints in a small region. Most of the parameter space will require a discrete symmetry in order to accommodate the $Z \rightarrow b_L \bar{b}_L$ constraint.

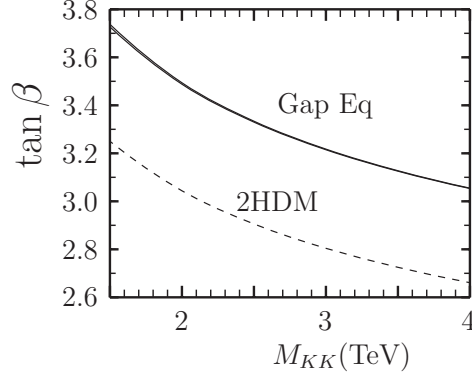


FIG. 7: The solutions for $\tan \beta$ v.s. M_{KK} in the 2HDM approach and from the gap equation.

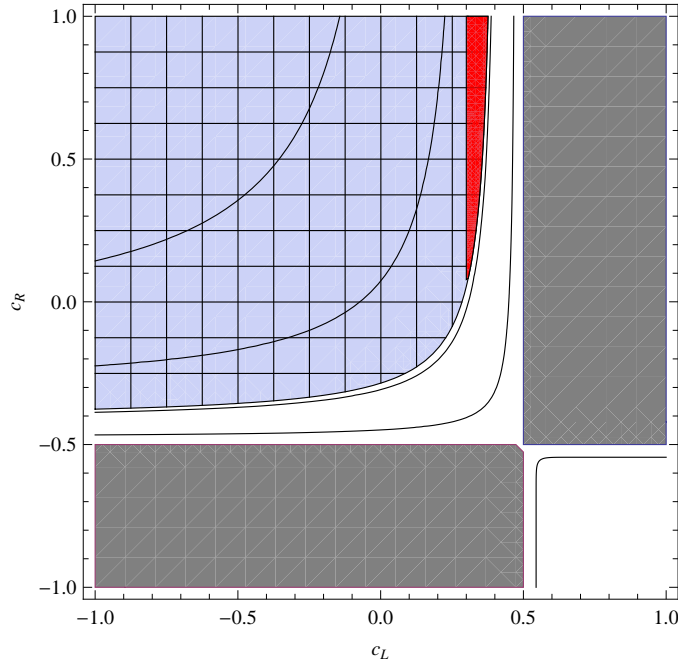


FIG. 8: Contours of g_t in c_L^3, c_R^3 plane for $M_{KK} = 1.5$ TeV. The solid lines are g_t contours for $g_t = \{4, 3, 2, 1, 0.125\}$ from top-left to bottom right. The dark regions are where g_L and g_R have opposite signs, and hence are inconsistent with the condensate scenario. The small red region gives a good fit to $Z \rightarrow b_L \bar{b}_L$, without the additional P_{LR} symmetry.

V. FLAVOR CHANGING NEUTRAL CURRENT EFFECTS.

Now we consider the full Yukawa sector for the 2HDM model, including light quarks. The Lagrangian is given by

$$\mathcal{L}_Y = \lambda_{ij}^d \bar{Q}_{Li} d_{jR} H + g_t \bar{Q}_{3L} t_R \tilde{\Phi} + \lambda_{ij}^u \bar{Q}_{iL} u_{jR} \tilde{H} + h.c. \quad (31)$$

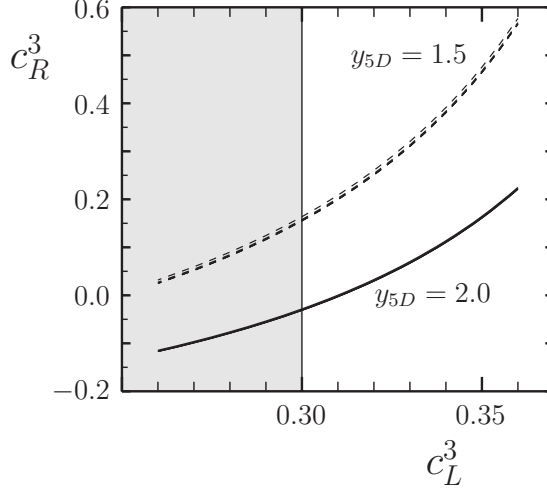


FIG. 9: The solution for bulk mass parameters c_L^3 and c_R^3 with two representative 5D Yukawa couplings. The KK mass is varied from 1.5 TeV to 4.0 TeV. The shaded areas are excluded by the $Z \rightarrow b_L \bar{b}_L$.

where i, j are family indices and other notations are standard. The condensate Φ only couples to the third family. After the rotation to remove scalar kinetic mixing given in Eq. (13), we get

$$\mathcal{L}_Y = \lambda_{ij}^d \overline{Q}_{Li} d_{jR} \hat{H} + (\lambda_{ij}^u - \lambda_{33}^u) \overline{Q}_{iL} u_{jR} \tilde{\hat{H}} + \frac{1}{\sqrt{\epsilon}} \overline{Q}_{3L} t_R \tilde{\hat{\Phi}} + h.c. \quad (32)$$

After spontaneous symmetry breaking, the mass matrices are given by replacing $\hat{H} \rightarrow v_H/\sqrt{2}$ and $\hat{\Phi} \rightarrow v_\phi/\sqrt{2}$. The down type mass matrix is proportional to λ_{ij}^d , so diagonalizing this also diagonalizes the Yukawa matrix. There are therefore no tree level flavor changing neutral currents for the d -type quarks.

However, the up type quarks receive masses from VEVs of both doublets. This implies that when we diagonalize the u -quark mass matrix, we will not simultaneously diagonalize the Yukawa matrix. This leads to tree level FCNC, which are obviously a matter of concern. While we defer a full study of these effects for later work, here we argue that they will not be intolerable. First, we can write the mass matrix for up-type quarks as

$$\mathcal{M}_{ij}^u = -\frac{1}{\sqrt{2}} \begin{pmatrix} \lambda_{11}^u v_H & \lambda_{12}^u v_H & \lambda_{13}^u v_H \\ \lambda_{21}^u v_H & \lambda_{22}^u v_H & \lambda_{23}^u v_H \\ \lambda_{31}^u v_H & \lambda_{32}^u v_H & \frac{1}{\sqrt{\epsilon}} v_\phi \end{pmatrix}. \quad (33)$$

This lets us rewrite the Yukawa sector (32) as

$$\mathcal{L}_Y = -\frac{\sqrt{2}\mathcal{M}_{ij}^d}{v\sin\beta}\overline{Q_{Li}}d_{jR}\hat{H} - \frac{\sqrt{2}\mathcal{M}_{ij}^u}{v\sin\beta}\overline{Q_{iL}}u_{jR}\tilde{\hat{H}} + \frac{1}{\sqrt{\epsilon}}\overline{Q_{3L}}t_R\left(\tilde{\hat{\Phi}} - \frac{\tilde{\hat{H}}}{\tan\beta}\right) + h.c. \quad (34)$$

It is manifestly clear that FCNCs arise *only* from the last term.

There are three effects suppressing light quark FCNCs from this term. The first two arise from the expansion of the Higgs doublets in terms of physical fields. Note that the combination $\sin\beta\tilde{\hat{\Phi}} - \cos\beta\tilde{\hat{H}}$ has no VeV, and so contains purely the *physical* charged and pseudoscalar fields. Because of the near equality of the mixing angles α and β , it also is almost pure H_0 . This means that flavour changing currents through the neutral scalars are either suppressed by the relatively large mass M_H^2 , if they proceed via H_0 ; or by the small parameter $\sin(\beta - \alpha)$, if they proceed via h_0 .

The third effect suppressing light quark FCNCs comes from the flavor structure of RS. The gauge basis we have used above is close to the 4D mass basis; the c -parameters of the 5D fermions naturally generate hierarchical Yukawa matrices. This implies that $\overline{Q_{3L}}$ and t_R states contain only small admixtures of the up and charm quarks, suppressing FCNCs that do not involve t-quarks. We expect that the combination of these factors suffices to push new contributions from this model to measured observables to sufficiently small values.

Finally, we briefly comment on the first two terms in Eq. (34). These terms involve the dominant coupling of the light scalar h_0 to fermions, since the last term introduces a heavily suppressed coupling as already noted. We note that the neutral sector of \hat{H} is $\sin\alpha h_0 + \cos\alpha H_0$ (see Eq. (25)). Then we see that the Yukawa matrix for h_0 is $-\sqrt{2}(\mathcal{M}_{ij}^d/v)\sin\alpha/\sin\beta$, as already claimed in section IV.

VI. CONCLUSIONS

We have derived an effective 2HDM from the Randall-Sundrum scenario by using a IR brane localized Higgs field and a second scalar doublet obtained from t-quark condensation induced by KK gluon exchange. The first KK gluon mode is dominant, and we take its mass M_{KK} to be the cutoff for the effective theory. The scalar kinetic mixing and mass mixing are calculated in the one fermion bubble approximation. After the diagonalization that gives canonical kinetic terms the resulting 2HDM is surprisingly simple. The model has no tree level FCNC effects for the d-type quarks and hence does not run afoul of constraints

from B meson and kaon decays. On the other hand we expect sizable FCNC decays of the t-quark which can be searched for at the LHC, and we defer this and other flavor physics issues to a later work. Moreover, one can show that the model does not support spontaneous T-violation.

We obtained the t-quark mass in two different ways. The first was to use the effective 2HDM, which after SSB yields the result that m_t is proportional to $\cos\beta$, where β is the mixing angle of the two VeV's. For the second approach we derive an inhomogeneous gap equation for m_t , using the fermion bubble approximation and the brane Higgs. The two results are in qualitative agreement, and quantitatively they differ by no more than 20%. This is well within the N_c^{-1} accuracy of our approximation.

As in any 2HDM, there are five physical spin-0 particles. The model predicts that the pair of charged scalars and the pseudoscalar are degenerate at tree level. They are expected to be close to M_{KK} , as is the heavier real scalar. The remaining scalar is at the electroweak scale, and has couplings to fermions and gauge bosons that are essentially the same as the SM Higgs. The heavier scalar has a small coupling to the electroweak gauge bosons, and hence can only be produced at the LHC via gluon fusion through quark loops.

Since the second Higgs owes its origin to the formation of a t-condensate due to KK gluons, it is not surprising that the properties of the resulting 2HDM depend on the localization c -parameters of the t-quark. We note that there is a small region of parameter space that satisfies the EWPT and the constraints of $Z \rightarrow b_L \bar{b}_L$ in the custodial $SU(2)_L \times SU(2)_R \times U(1)_X$ RS model without an additional discrete symmetry.

Acknowledgments

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